

Bimodal energy relaxation in quasi-one-dimensional compounds with a commensurate modulated ground state

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We show that the low temperature ($T < 0.5$ K) time dependent non-exponential energy relaxation of quasi-one-dimensional (quasi-1D) compounds strongly differ according to the nature of their modulated ground state. For incommensurate ground states, such as in $(\text{TMTSF})_2\text{PF}_6$ the relaxation time distribution is homogeneously shifted to larger time when the duration of the heat input is increased, and exhibits in addition a scaling between the width and the position of the peak in the relaxation time distribution, $w^2 \sim \ln(\tau_m)$. For a commensurate ground state, as in $(\text{TMTTF})_2\text{PF}_6$, the relaxation time spectra show a bimodal character with a weight transfer between well separated slow and fast entities. Our interpretation is based on the dynamics of defects in the modulated structure, which depend crucially on the degree of commensurability.

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Non exponential relaxation is the fingerprint of disorder in a system subject to an external perturbation [1]. It is the consequence of a broad distribution of relaxation times. The question of determining the nature and spatial extent of correlated objects responsible for slow dynamics in spin glasses and glasses, has been debated for several decades. In spin glasses, the memory and chaos experiments [2] give information about excitations in real space. In glasses there is a consensus that structural relaxation occurs on correlated domains of a nanometer size [3]. In colloidal suspensions, a direct three dimensional (3D) imaging of the cooperative motion in structural relaxation has been obtained on micrometer scales [4]. A general framework for describing spatial and temporal heterogeneities in aging systems has been proposed in Ref. [5]. A new class of mesoscopic glass, with characteristic scales of micrometers, has been recognized among quasi-one-dimensional (quasi 1D) systems [6]. Heat relaxation at very low temperature T , below 0.5 K, in these systems, has been shown to be non exponential and to exhibit aging [6] in which the entire wide spectrum of relaxation times is shifted to longer times with increasing the duration of the very small heat perturbation. We suggest here on the basis of heat relaxation experiments that the formation of correlated objects in modulated quasi-1D systems, is qualitatively different in commensurate and incommensurate systems.

Quasi 1D systems exhibit a variety of modulated electronic ground states at low-temperature both in the charge channel, such as charge density wave (CDW), and in the spin channel, such as spin density wave (SDW), spin-Peierls (SP) or antiferromagnetic ground state. Extensive investigations of the low- T thermodynamical properties [6, 7, 8, 9] have shown that all of these systems demonstrate well-known universal features of glasses regardless the actual ground state. This in-

volves an extra contribution to the specific heat C_p coming from low energy excitations, and non exponential heat relaxation dynamics for $T < 0.5$ K (T -range dominated by these excitations). “Aging” effects are also involved, where heat relaxation and correspondingly C_p depend on the duration of a small heat perturbation (a few mK) [6, 7, 8]. Extremely long relaxation times (up to 10^4 s) show that these excitations are only weakly coupled to the phonon heat-bath [8], contrary to conventional glasses.

The origin of these excitations can be inferred from the low frequency dynamics of collective excitations in the charge channel. Low frequency dielectric response of CDW [10], SDW [11] which also has CDW component [12], all having a modulated charge density, is characterized by an overdamped (relaxational) mode slowing down with decreasing temperature. This has been recognized in the case of CDW and SDW (DW) as the sign of a glass transition at the level of electronic superstructure. Freezing is naturally generated by screening effects and occurs if there are not enough free carriers to screen efficiently elastic DW phase deformations. Consequently, the characteristic scale of a DW glass is given by the phase coherence length, which is of the order of a μm . Below the glass transition temperature T_g the low frequency dynamics and therefore the low energy excitations are governed by the remaining degrees of freedom – topological or plastic deformations such as solitons, domain walls or dislocation loops.

In this Letter we contrast two isostructural systems, $(\text{TMTSF})_2\text{PF}_6$, which has an incommensurate mixed SDW-CDW ground state below 12 K, and $(\text{TMTTF})_2\text{PF}_6$, which exhibits a charge order transition with a ferroelectric behavior at 70 K [13], and a spin-Peierls transition at 12 K, both being commensurate to the underlying lattice (in the following we use

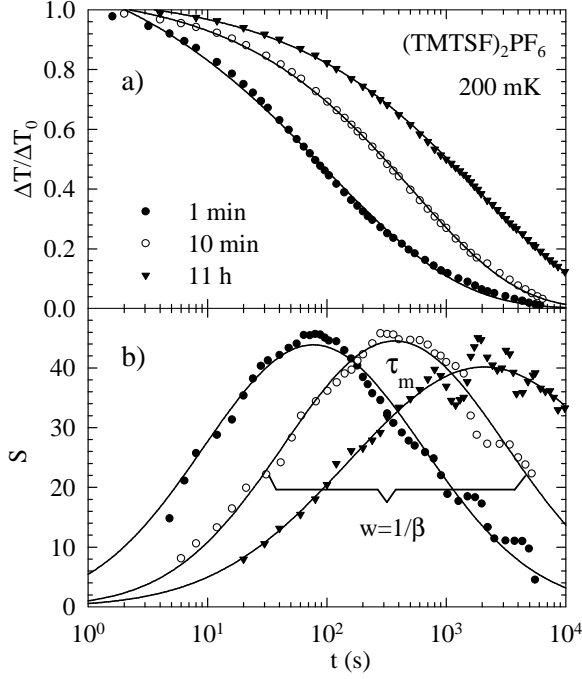


FIG. 1: *Homogeneous* broadening and shift of the heat relaxation with increasing t_p (aging) for SDW-PF₆. a) Relaxation $\Delta T / \Delta T_0$ vs. $\log t$ with corresponding stretched exponential fit $\exp(-t/\tau)^\beta$, b) $S(t) = d(\Delta T / \Delta T_0) / d(\ln t)$ fitted to a Gaussian function $G(\ln t)$ with a width w . The direct estimate of width at half height shown in b) is 2.17 decades in agreement with $1/\beta = 2.15$ and $w = 2.1$.

the notation SDW-PF₆ and SP-PF₆ respectively). We will show how the degree of commensurability affects the very low- T dynamics and discuss these findings in view of the properties of solitonic excitations.

We have measured the low- T heat capacity by standard relaxation method. The system is first stabilized at a temperature T_0 . Then a small power is fed to the sample during a pumping time t_p . The temperature is increased to the new, slightly higher value $T_0 + \Delta T$ ($\Delta T / T_0 < 0.05$). At $t = t_p$ the power is switched off and heat relaxation from the sample to the cold source towards equilibrium at T_0 is recorded as a function of time. For all the CDW and SDW systems studied previously, below 1 K relaxation starts to deviate from exponential on lowering T and also exhibits a strong dependence on t_p [6, 7] (Fig. 1). A better characterization of non exponential relaxation can be obtained by the relaxation rate $S(t) = d(\Delta T / \Delta T_0) / d(\ln t)$. $\Delta T(t) / \Delta T_0$ can be written in terms of a distribution of relaxation times $P(\ln \tau)$ as

$$\frac{\Delta T(t)}{\Delta T_0} = \int_{\ln \tau_0}^{\infty} P(\ln \tau) \exp(-t/\tau) d \ln \tau, \quad (1)$$

where τ_0 is the microscopic time. Since $P(\ln \tau)$ varies slowly with τ we replace the exponential by a step function, from what we deduce $S(t) = P(\ln t)$ [14]. Like in

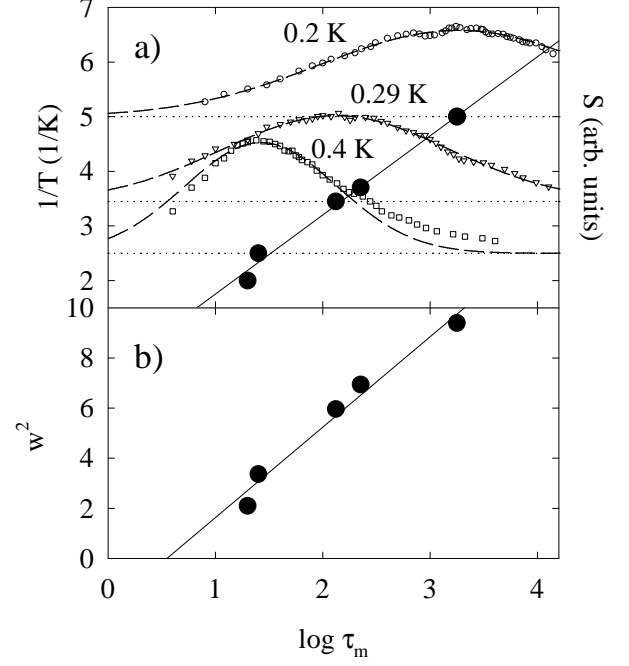


FIG. 2: a) An “inverted” Arrhenius plot, $1/T$ vs. $\log \tau_m$, obtained for relaxations in thermodynamical equilibrium ($\tau_m = \tau_{eq}$) of SDW-PF₆. Relaxation rates and their fits are only shown for three selected temperatures among five. b) The $w^2 \sim \ln \tau_m$ scaling is well obeyed as expected from central limit theorem. The activation energy is $E_a = 1.7$ K.

spin glasses there is a “homogeneous” broadening and shifting of the relaxation peak with t_p (see Fig. 1-(b)). However, since our system eventually reaches thermodynamic equilibrium at $T_0 + \Delta T$ for sufficiently long t_p [7], aging is then interrupted [15].

Non exponential relaxation can be fitted by a “stretched exponential” $\Delta T / \Delta T_0 \sim \exp[-(t/\tau)^\beta]$ [6], while the relaxation rate $S(t)$ can be fitted by a Gaussian distribution in $\ln t$: $G(\ln t) = (1/\sqrt{2\pi w}) \exp(-\ln^2(t/\tau_m)/2w^2)$, centered on $\ln \tau_m$ and with a width w . Indeed, the logarithmic derivative of a stretched exponential $d(\exp(-(t/\tau)^\beta))/d(\ln t)$ is almost a Gaussian with $\tau_m \sim \tau$ and $1/w \sim \beta$, as shown in Fig. 1. The same w can also be extracted directly from $S(t)$ curves as the width at half height.

While the stretched exponential has no known justification other than phenomenological [1, 3], Castaing and Souletie [16] argued that a Gaussian distribution in $\ln t$ can be justified in a framework that associates a hierarchical criterion and a usual scaling approach in the renormalization of distances. Similarly to turbulence they use a cascade of events to relate the macroscopic and microscopic scales. After iterating the renormalization group dynamical scaling takes the form $\tau/\tau_0 \sim (\xi/\xi_0)^z$, where z is the dynamical exponent. The time scale distributions are log-normal: $P(\ln \tau) \equiv G(\ln t)$ is Gaussian.

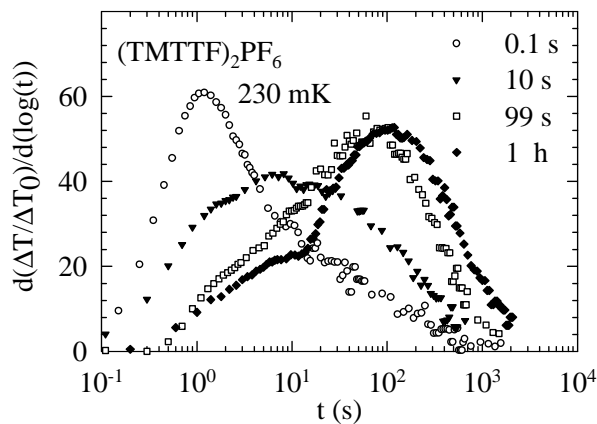


FIG. 3: *Bimodal* relaxation with redistribution of the spectral weight and saturation for $t_{eq} \sim 10^2$ s for the spin-Peierls compound SP-PF₆. The values of t_p are indicated on the figure

The variance w^2 scales with the most probable value $\ln(t/\tau_m)$ [17] following the central limit theorem. It is demonstrated in Fig. 2 for SDW-PF₆ in thermodynamical equilibrium [9] for three temperatures chosen below 0.5 K. This yields immediately the temperature dependence of the stretched exponential exponent $\beta(T) \sim \sqrt{T}$, as $\beta \sim 1/w \sim 1/\sqrt{\ln \tau_m}$ (Fig. 2b), so that $\ln \tau_m \sim 1/T$ indicates a thermally activated process with activation energy $E_a = 1.7$ K. At this point we do not intend to go deeper into the different models related to $\beta(T)$ for $T \rightarrow 0$ (see ref. 6 and references therein). We stress that the $w^2 \sim \ln \tau_m$ scaling is quite convincing in the range of interest, and also for TaS₃, a prototype CDW system, for which as well as for SDW-PF₆ a glass transition has been found at higher T [10].

However, in the case of SP-PF₆ [9] heat relaxation, while still non exponential, exhibits quite different properties. Instead of the homogeneous broadening and shift of the spectra of SDW-PF₆ as in a real “aging” effects, SP-PF₆ demonstrates “discrete bands” of relaxation times. Their distributions are narrower and the relaxation times are smaller than in SDW-PF₆ (i.e. saturation of aging is reached within a few tens of minutes as compared to days in first system at a similar T). Fig. 3 shows a bimodal redistribution of the relaxation spectrum indicating that heat perturbation modifies different parts of the spectrum in contrast to the homogeneous shift of the entire spectrum as in Fig. 1.

Even in thermodynamical equilibrium the $S(t)$ curves are not symmetric, indicating at least two different relaxation processes. However the total w (Fig. 4) can be obtained, but the variation $w^2 \sim \ln \tau_m$ is very peculiar. We have successfully deconvoluted $S(t)$ in two dynamically distinct Gaussian contributions, represented by dashed lines in Fig. 4a. At higher T , the fast contribution plays the dominant role in the dynamics, but the

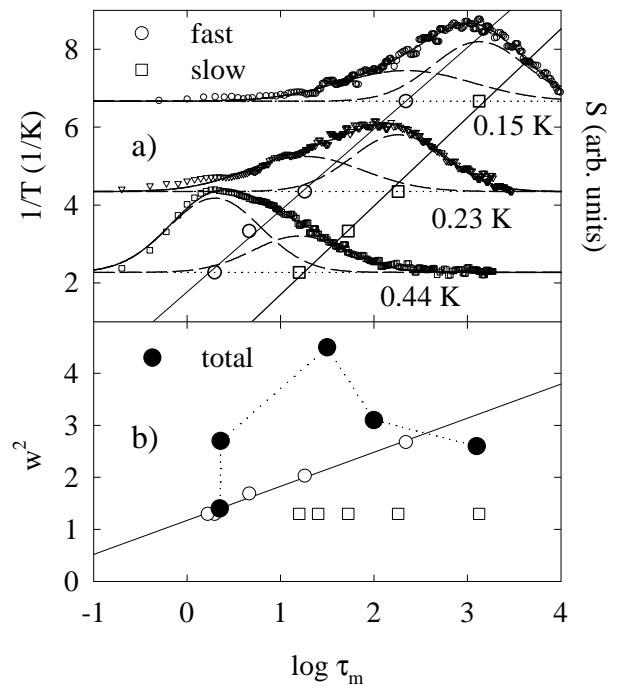


FIG. 4: a) An “inverted” Arrhenius plot, $1/T$ vs. $\log \tau_m$, obtained for the relaxations in thermodynamical equilibrium of the spin-Peierls compound SP-PF₆, together with a representation of corresponding relaxation rates for slow and fast processes and their fits, with $E_{a,s} = E_{a,f} = 1.0$ K b) The $w^2 \sim \ln \tau_m$ scaling does not work for the total relaxation, however it might work for each deconvoluted process.

relative weight of slow modes in the dynamics grows on decreasing the temperature. A similar behavior occurs at a fixed temperature if we change the duration t_p of the heat perturbation (Fig. 3). There is an evident weight transfer from one to the other subsystem as thermodynamic equilibrium is approached. However, without the fast entities, there would be almost no response within 10^2 - 10^3 seconds at the lowest T . The distribution of fast modes is pushed to 10^3 seconds at 85 mK, and the apparent activated behavior (Fig. 4a) turns into a saturation indicating possible tunneling effects (not included in Fig. 4).

There are drastic consequences on the proper definition and determination of the specific heat C_p in these systems as it becomes strongly nonlinear [8]. From the deconvoluted entities in thermodynamic equilibrium (Fig. 4a), it is possible to calculate the corresponding C_p , fast and slow. Their participation changes progressively from approximately 100 % for the fast modes at ~ 650 mK to almost 90% for the slow modes below 400 mK.

It appears that the incommensurate CDW compound o-TaS₃ shares all its properties related to the shift of the dynamics with SDW-PF₆. Both are nominally incommensurate systems. On the other hand, (TMTTF)₂Br shares its multimodal dynamics proper-

ties with SP-PF₆ [9]. (TMTTF)₂Br being the brother compound of SP-PF₆ is a commensurate antiferromagnet. We conclude our Letter by proposing an explanation to these distinctly different dynamics. The models generalizing Refs.[18, 19, 20] are investigated in more details in Ref. [21].

As proposed by Larkin [18] and Ovchinnikov [19], the local deformations of the DW in the strong pinning limit are bisolitons, obtained by minimizing the elastic and pinning energies, and characterized by a ground state at energy E_0 , separated from another state at energy $E_0 + \Delta E$ by a “bounce” state at energy $E_0 + \Delta V$, therefore defining an effective two-level system. As a qualitative difference between commensurate and incommensurate systems[21], one has $\Delta E = 0$ in the commensurate case [22] and $\Delta E > 0$ in the incommensurate case, so that the populations of the effective two level system does not couple to temperature variations in the commensurate case. We suggest here that, in addition, the degeneracy of the classical ground state in the commensurate case is lifted by quantum fluctuations that restore a finite heat response. The long time dynamics in the incommensurate case is due to collective effects, as suggested by the $w^2 \sim \ln \tau_m$ scaling on Fig. 1, in agreement with Ref. [21]. The issue of the long time dynamics in the commensurate case, involving both quantum and collective effects, is mainly an open theoretical question.

To explain why the spectrum of relaxation time is bimodal in the commensurate case we note that experimental data suggest that in the commensurate system there exist two entities that relax with approximately the same energy barrier but with values of the “microscopic time” differing by one order of magnitude. Possible candidates for these two entities would be (i) the dipole solitonic excitations generated by the strong pinning impurities; and (ii) self-induced disorder due to 2π -solitons between micro-domains separating the two spin-Peierls ground states. Alternatively one may observe that SP-PF₆ has a charge gap $\Delta_C \simeq 200$ K coexisting with a spin gap $\Delta_S \simeq 20$ K, deduced from the spin-Peierls transition temperature from the mean field approximation. The coherence length in the charge channel $\xi_C = \hbar v_F / \Delta_C \simeq 10a_0$ is thus much smaller than the coherence length in the spin channel $\xi_S = \hbar v_F / \Delta_S \simeq 100a_0$ (where a_0 is the lattice parameter and v_F the Fermi velocity). The two channels are weakly coupled so that the charge soliton, being smaller, relaxes faster than the spin soliton. The pinning potential and thus the energy barriers are identical in the two sectors so that the activation energies are identical but the microscopic times are different, in agreement with Fig. 4.

In conclusion we have shown manifestations of correlated objects interpreted as the collective dynamics of a disordered soliton lattice inherent to the modulated electronic ground state of DW. We demonstrated that

a simple scaling feature of the parameters of the relaxation time distribution is well obeyed for the homogeneously broadened response for incommensurate systems. By contrast the heat response of commensurate systems is bimodal. We stress that our results should be also compared to other experiments on the dynamical heterogeneity in the relaxations of supercooled liquids or spin glasses [24, 25, 26], and to the case of bimodal relaxation [23] observed in a fragile glass former at the approach of a glass transition. However, the microscopic picture is different, in spite of a similar phenomenology.

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